

Remote Inflation as hybrid-like sneutrino/MSSM inflation

Tomohiro Matsuda¹

*Laboratory of Physics, Saitama Institute of Technology,
Fusaiji, Okabe-machi, Saitama 369-0293, Japan*

Abstract

A new scenario of hybrid-like inflation is considered for sneutrino and MSSM fields. Contrary to the usual hybrid inflation model, the direct coupling between a trigger field and the sneutrino/MSSM inflaton field is not necessary for the scenario. The dissipation and the radiation from the sneutrino/MSSM inflaton can be written explicitly by using the Yukawa couplings. Remote inflation does not require the shift symmetry or cancellation in solving the η problem.

¹matsuda@sit.ac.jp

1 Introduction

According to the recent cosmological observations, inflation has become the paradigm for early cosmology and appears to be the most successful cosmological model giving the primordial large scale structure of the Universe. Considering typical slow-roll inflation, there are at least two scenarios for the velocity damping: the original (supercooled) inflation and warm inflation in which a strong diffusion produces significant friction for the inflaton motion[1]. Recently, hybrid-like extension of warm inflation has been proposed in Ref.[2], in which a new possibility is shown that the radiation produced concurrently during warm inflation may keep symmetry restoration in a remote sector. As a consequence, the false-vacuum energy of the remote sector may dominate the energy density during inflation. The situation is similar to thermal inflation, however in contrast to the standard thermal inflation model, the radiation in remote inflation is sourced by the dissipation of the inflaton field. Remote inflation is a new inflationary model based on thermal inflation[3], hybrid inflation[4] and warm inflation[1] and is expected to solve problems in these classic models. In this paper we consider remote inflation for sneutrino inflation - inflation caused by the scalar supersymmetric partner of a heavy singlet neutrino[5, 44] - and MSSM inflation - inflation caused by the fields in the minimal supersymmetric standard model (MSSM)[6, 7].

Remote inflation is based on warm inflation. Warm inflation for a sneutrino field has been discussed in Ref.[8], in which monomial potentials constructed with only the right-handed sneutrino field realize warm inflationary regime for chaotic inflation. Since the conditions for warm inflation and supercooled inflation are exclusive, the parameter space for successful sneutrino inflation must be complemented by considering warm inflation. With regard to the problems and the benefits of sneutrino/MSSM, hybrid and warm inflation scenarios, remote inflation may give a very attractive situation in which hybrid-like inflation is realized in the warm inflationary regime without introducing additional direct coupling between the inflaton and a trigger field. Applying remote inflation to sneutrino and MSSM fields, it may be possible to find a successful scenario in which the baryon asymmetry of the Universe is created by the non-equilibrium decay of the

inflaton.²

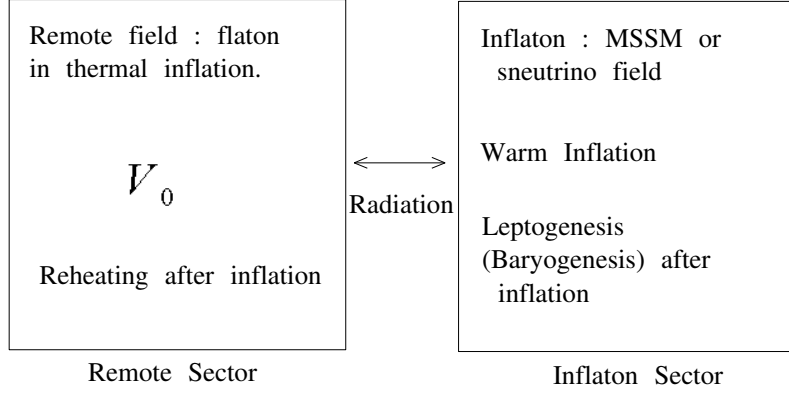


Figure 1: Remote inflation resembles hybrid inflation. Remote field plays the role of the trigger field in hybrid inflation, which causes reheating the Universe. The explicit interaction in the conventional hybrid potential is replaced by the thermal background. Also, the remote sector resembles thermal inflation. The source of the radiation is the dissipation in the inflaton sector, where warm inflation occurs.

The idea of remote inflation is very simple. Considering the standard hybrid inflation model, the typical form of the hybrid-type potential is given by

$$V(\phi, \sigma) = V(\phi) + \frac{g^2}{2} \sigma^2 \phi^2 + \frac{\lambda}{4} (\sigma^2 - M^2)^2, \quad (1.1)$$

where the vacuum energy during inflation is $V_0 \simeq \lambda M^4/4$. The trigger field σ stays in its false-vacuum state defined by $\sigma = 0$ for $\phi > \phi_c$. Here the critical point that defines the end of inflation is given by

$$\phi_c \equiv \sqrt{\frac{\lambda}{g}} M. \quad (1.2)$$

Contrary to the usual hybrid inflation scenario, either normal, inverted [9] brane-motivated [10, 11, 12], or thermal [13], remote inflation does not have direct couplings between inflaton and trigger fields. The trigger fields σ_i in the remote inflation model are placed in a remote sector and stay in their false vacuum states during inflation due to the thermal effects created by the diffusion of the inflaton field. Namely, the force that keeps

²See Fig.1.

σ_i at their false-vacuum state is not the effective mass $m_\sigma^2 \sim g^2\phi^2$ caused by the direct coupling between the inflaton and the trigger field, but the thermal effect of the radiation concurrently produced during warm inflation.

Remote inflation requires symmetry restoration in the remote sector. To illustrate some typical features of finite temperature effects in the remote sector, here we consider a real scalar field and a potential:

$$\begin{aligned}\mathcal{L} &= \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma - V(\sigma) \\ V(\sigma) &= V_0 - \frac{1}{2}m_\sigma^2\sigma^2 + \frac{1}{4}\lambda\sigma^4,\end{aligned}\tag{1.3}$$

where V_0 is tuned so that the cosmological constant vanishes at the true minimum. The phenomenon of high-temperature symmetry restoration can be understood by the finite-temperature effective potential given by [14]

$$V_T(\sigma_c) = V(\sigma_c) + \frac{T^4}{2\pi^2} \int_0^\infty dx \ln \left[1 - \exp \left(-\sqrt{x^2 + \frac{-m_\sigma^2 + 3\lambda\sigma_c^2}{T^2}} \right) \right], \tag{1.4}$$

where $V(\sigma_c)$ is the one-loop potential for zero-temperature with the classical field σ_c :

$$V(\sigma_c) = -\frac{1}{2}m_\sigma^2\sigma_c^2 + \frac{1}{4}\lambda\sigma_c^4 + \frac{1}{64\pi^2} (-m_\sigma^2 + 3\lambda\sigma_c^2)^2 \ln \left(\frac{-m_\sigma^2 + 3\lambda\sigma_c^2}{\mu^2} \right), \tag{1.5}$$

where μ is a renormalization mass scale. At high temperatures, V_T can be expanded near the symmetric point ($\sigma_c = 0$) as

$$V_T \simeq V(\sigma_c) + \frac{1}{8}\lambda T^2\sigma_c^2 + \mathcal{O}(T^4), \tag{1.6}$$

which suggests that the temperature-corrected effective mass at $\sigma_c = 0$ changes sign at the critical temperature given by

$$T_c \simeq \frac{2m_\sigma}{\lambda^{1/2}}. \tag{1.7}$$

More generally, a background thermal bath which couples to σ can be introduced. Then we obtain a typical form of the potential with a thermal correction term, which is given by

$$V = V_0 + \left(g^2 T^2 - \frac{1}{2}m_\sigma^2 \right) \sigma^2 + \dots, \tag{1.8}$$

where g denotes the effective coupling of σ to the fields in the thermal bath. In this case, the critical temperature is given by

$$T_c \simeq \frac{m_\sigma}{2g}. \tag{1.9}$$

In this paper, we assume that the couplings to the background fields are significant and T_c is given by $T_c \simeq \frac{m_\sigma}{2g} \sim m_\sigma$. Here $\lambda \ll 1$ is possible for a flat potential in supersymmetric theory, which gives typically the false-vacuum energy $V_0 \sim m_\sigma^2 M^2$ for the quartic potential, where $M^2 \sim m_\sigma^2/\lambda \gg m_\sigma^2$ denotes the vacuum expectation value of the field σ in the true vacuum.

2 Remote inflation with dissipative sneutrino

The interaction of the sneutrino inflaton field with other fields is very important in sneutrino inflation, since Yukawa couplings are always needed for the seesaw mechanism[15] that gives the tiny masses to the light neutrinos. The heavy singlet neutrinos usually have masses 10^{10} to 10^{15} GeV, but one may consider a smaller mass for a singlet neutrino. For our scenario the simplest and sufficient model is only one sneutrino acting as the inflaton whose dissipation sources radiation during inflation. The interaction relevant for the model is extracted from the supersymmetric version of the seesaw model as

$$\mathcal{L}_I \sim -|h_N|^2 |\phi_N|^2 |H_u|^2 + h_t H_u \bar{t}_R t_L + h.c., \quad (2.1)$$

where ϕ_N is the sneutrino inflaton, H_u is the Higgs doublet giving masses to the up quarks and t_R, t_L are the top quarks. Dissipative effects are caused by the coherent excitation of the H_u field, which then decays into fermions (standard-model quarks) with the decay rate

$$\Gamma_t = \frac{3h_t^2}{16\pi} m_{H_u} \sim \frac{3h_t^2}{16\pi} h_N \phi_N. \quad (2.2)$$

The condition for the thermalization of the decay product is given by

$$\Gamma_t > H, \quad (2.3)$$

which puts a natural upper bound for the energy scale of warm inflation. Following Ref.[1] and [8], the typical form of the dissipative coefficient Γ for warm inflation is given by

$$\Gamma \simeq \frac{\sqrt{\pi}}{20} \frac{h_N^2 h_t^2}{(4\pi)^2} \phi_N \equiv A_N \phi_N. \quad (2.4)$$

The strength of the thermal damping is measured by the rate r_Γ given by

$$r_\Gamma \equiv \frac{\Gamma}{3H}, \quad (2.5)$$

which can be used to rewrite the field equation as

$$\ddot{\phi} + 3H(1 + r_\Gamma)\dot{\phi}_N + V(\phi_N, T)_{\phi_N} = 0, \quad (2.6)$$

where the subscript denotes the derivative with respect to the field. Considering a typical quadratic potential for the singlet neutrino, a simple form can be used for the potential;

$$V(\phi_N, T)_{\phi_N} \simeq M_N^2 \phi_N, \quad (2.7)$$

where M_N is the mass of the sneutrino. From the above equations the slow-roll velocity of the sneutrino inflaton for the strongly dissipating regime ($r_\Gamma > 1$) is given by

$$\dot{\phi}_N \simeq \frac{M_N^2 \phi_N}{\Gamma} \simeq \frac{M_N^2}{A_N}, \quad (2.8)$$

which does not depend on the temperature. Note that the slow-roll velocity is independent of the temperature of the Universe. The conventional slow-roll parameters are given by

$$\begin{aligned} \epsilon &\equiv \frac{M_p^2}{2} \left(\frac{V_\phi}{V} \right)^2 < (1 + r_\Gamma) \\ \eta &\equiv M_p^2 \left(\frac{V_{\phi\phi}}{V} \right) < (1 + r_\Gamma). \end{aligned} \quad (2.9)$$

In addition to the standard slow-roll conditions for the inflaton motion, slow-roll approximation for warm inflation requires slow-variation of Γ , which leads to an additional slow-roll parameter

$$\beta \equiv M_p^2 \left(\frac{\Gamma_{\phi_N} V_{\phi_N}}{\Gamma V} \right) \simeq \frac{M_N^2}{H^2} \sim \eta < (1 + r_\Gamma). \quad (2.10)$$

This equation gives a trivial condition for the present model. Similarly, T -dependence of Γ leads to a condition for the radiation[1]

$$c \equiv \left| \frac{d \ln \Gamma}{d \ln T} \right| = \left| \frac{d \ln \Gamma}{d \ln \phi_N} \frac{d \ln \phi_N}{d \ln T} \right| = 4, \quad (2.11)$$

which does not satisfy the required condition $c < 4$. Here $c = 4$ is the direct consequence of the quadratic potential, namely $V \propto \phi_N^2$ with $n = 2$. However, in more general situations the effective potential may have other higher terms with $n > 2$, although they may be small compared with the quadratic term. Therefore, the effective n would be $n = 2 + \epsilon_N > 2$, where $\epsilon_N \ll 1$. A small shift from $n = 2$ is trivial and it can be

disregarded in other scenarios, but in hybrid-type warm inflation the shift is crucial for warm inflation[16]. Moreover, if initially the scalar field gets a value as large as the Planck (or the fundamental) scale, the potential at that time may have significant higher corrections with large n . Therefore, in this paper we consider the quadratic potential, but the effective value of n is slightly larger than $n = 2$. Other possibilities, the potential characterized by higher terms with $n > 2$, or a flat potential characterized by $\sim \log \phi$ are discussed in section 2.4.

During inflation, the energy conservation of the radiation energy density gives the equation

$$\dot{\rho}_r + 4H\rho_r = \Gamma\dot{\phi}^2, \quad (2.12)$$

which leads to the temperature during inflation given by

$$T^4 = \frac{45}{2\pi^2 g_*} r_\Gamma \dot{\phi}^2, \quad (2.13)$$

where $g_* \sim \mathcal{O}(10^2)$ is the effective number of the light degrees of freedom at the temperature. A useful expression can be found from the above equation, which is given by

$$T^4 \simeq M_N^4 \frac{\phi_N}{A_N H}, \quad (2.14)$$

or equivalently;

$$r_T^4 \equiv \left(\frac{T}{H}\right)^4 \simeq \eta^2 \frac{\phi_N}{A_N H}. \quad (2.15)$$

Since the inflationary expansion is not viable when the radiation dominates the Universe, we also consider the obvious condition for the radiation density ρ_r , which is given by

$$\rho_r < V_0. \quad (2.16)$$

In contrast with standard sneutrino-warm inflation[8], in which the above condition determines the end of inflation, this equation leads to the mild condition for remote inflation;

$$\phi_N < \phi_N^{(r)} \equiv A_N H \left(\frac{V_0}{M_N^4} \right). \quad (2.17)$$

If the inflaton reaches the breaking-point of the slow-roll condition ($\phi_N^{(e)}$) ahead of the critical temperature ($T = T_c$), the temperature at the end of inflation ($T = T_e$) is still higher than the critical temperature ($T_e > T_c$). The number of e-foldings is given by

$$N_e \simeq - \int_{\phi_N^{(e)}}^{\phi_N^{(i)}} \frac{3H^2}{\dot{\phi}} d\phi + N_{th} \simeq \frac{1 + r_\Gamma}{\eta} \ln \frac{\phi_N^{(i)}}{\phi_N^{(e)}} + N_{th}, \quad (2.18)$$

where the additional expansion of the Universe, denoted by N_{th} , is caused by thermal inflation that starts at $T = T_e$ and ends at $T = T_c$. Thermal inflation is defined for an inflationary expansion due to the vacuum energy of a false-vacuum state, but in contrast to remote inflation there is no significant source that supplies radiation during this epoch.³ The e-foldings of the expansion due to thermal inflation is given by $N_{th} \sim \ln(T_e/T_c)$. On the other hand, if inflation ends when the temperature reaches $T = T_c$ ahead of $\phi = \phi_N^{(e)}$, symmetry breaking in the remote sector occurs just at the end of remote inflation. In the latter case there is no thermal inflation after remote inflation [2].

2.1 Slow-roll conditions and the sources of the cosmological perturbations

In standard warm inflation, the effective potential depends on the temperature T , which can be expressed as $V(\phi_N, T)$. The damping rate Γ of the inflaton field may also depend on T and ϕ_i . Therefore, the trajectory of warm inflation is generally given by the inflaton and the temperature. Contrary to the standard warm inflation scenario, Γ in sneutrino inflation is independent of T . We thus find immediately that the damping of the inflaton ϕ_N is determined independent of the radiation. In contrast to the above argument, which suggests that the temperature is not important for the inflaton dynamics of dissipating (not necessarily warm) sneutrino inflation, the temperature is important for determining the end of inflation. As will be shown in this section, in sneutrino-remote inflation the phase transition in the remote sector triggers reheating and determines the end of inflation. Therefore, there are two sources of the curvature perturbations in remote inflation; the curvature perturbations created near the horizon exit and the one created at the inhomogeneous phase transition. In this paper we consider these two sources and compare the magnitude of the curvature perturbations.⁴ Although the conditions for remote inflation

³Breaking of the slow-roll condition does not always mean the end of dissipation. Here we assumed for simplicity that $\Gamma \propto \phi_N$ decreases rapidly after the critical point at $\phi_N = \phi_N^{(e)}$, and as a consequence, the dissipation caused by the inflaton oscillation is small during this period. Numerical studies are required in order to understand more details of the final stage of remote inflation, where oscillating inflaton may source radiation by dissipation, evaporation or decay.

⁴In addition to these sources, inhomogeneities of the diffusion rate $\delta\Gamma$ may be important and it may lead to a significant non-gaussian spectrum[18]. $\delta\Gamma$ causes inhomogeneities of the inflaton velocity, which

are highly model-dependent, we can find simple result at least for sneutrino inflation. Here the slow-roll conditions are defined by $\epsilon < r_{\Gamma}(\text{SDWI})$ and $\epsilon < 1(\text{WDWI})$, which determines $\phi_N^{(e)}$ as

$$\phi_N < \phi_N^{(e)} \simeq A_N \frac{H^3 M_p^2}{M_N^4} \sim A_N M_p \times \left(\frac{M_p}{M_N} \right) \eta^{-3/2} \quad (\text{SDWI}) \quad (2.20)$$

$$\phi_N < \phi_N^{(e)} \simeq M_p \eta^{-1} \quad (\text{WDWI}). \quad (2.21)$$

Sneutrino-remote inflation uses SDWI conditions. Note that the slow-roll condition is violated in the **outer region** of the field space, which becomes trivial for $\eta < 1$. Considering the critical temperature T_c , remote inflation is possible for the region given by

$$\phi > \phi_N^{(c)} \simeq A_N \frac{H T_c^4}{M_N^4} \simeq A_N \frac{H}{\eta^2} \left(\frac{T_c}{H} \right)^4 \simeq A_N \eta^{-2} r_{T_c}^3 \times T_c, \quad (2.22)$$

where r_{T_c} denotes the ratio $r_T \equiv T/H$ at the critical temperature. Considering the ratio between $\phi_N^{(e)}$ and $\phi_N^{(c)}$, we find

$$R_{e/c} \equiv \frac{\phi_N^{(e)}}{\phi_N^{(c)}} \simeq \frac{H^2 M_p^2}{T_c^4} \sim \frac{V_0}{T_c^4} \gg 1. \quad (2.23)$$

The equation shows that for sneutrino inflaton field, remote inflation is possible during $\phi_N^{(c)} < \phi < \phi_N^{(e)}$. Namely, we find in this model the phase transition in the remote sector determines the end of inflation.⁵

The spectrum of the cosmological perturbations created during inflation has been expected to be scale-invariant and Gaussian, but recent observations may suggest small anomalies. A small departure from exact scale-invariance and a certain non-Gaussian character [19, 20] can help reveal the dynamics of the inflation model. In particular, a tiny shift in the spectrum index $n - 1 \neq 0$ is a useful example[21] of scale invariance violation. The spectral index is used to determine the form of the inflaton potential.

 leads to inhomogeneities of the number of e-foldings[22]. Namely, considering the explicit form of the number of e-foldings given by Eq. (2.18), the perturbations δN [23] caused by $\delta \Gamma$ is given by

$$\delta N \sim N_e \frac{\delta \Gamma}{\Gamma} \sim N_e \left(\frac{\delta A_N}{A_N} + \frac{\delta \phi_N}{\phi_N} \right), \quad (2.19)$$

which may be as large as the conventional curvature perturbations. See Ref.[18] for more details on this topic.

⁵In contrast with remote inflation, the end of standard sneutrino-warm inflation is determined by the condition $\rho_r < V \sim M_N^2 \phi_N^2$ [8].

Observation of non-Gaussian character in the spectrum may lead to a significant bound for the mechanism of generating curvature perturbations. In order to explain how these anomalies are explained, many models of inflation and the generating mechanism of the curvature perturbations are proposed, in which the time when the perturbations are generated is very important. With regard to the strongly dissipating warm inflation, the curvature perturbations generated at the horizon exit are given by

$$\mathcal{R}_{warm}^{(ini)} = -H \frac{\delta q}{\rho + P} = H \frac{\dot{\phi} \delta \phi}{\dot{\phi}^2 + T_s} \simeq H \frac{\dot{\phi} \delta \phi}{T_s}, \quad (2.24)$$

which is not consistent with the expectation that the number of e-foldings in warm inflation is determined exclusively by the inflaton field ϕ . If the expectation is correct, the perturbation δN should be given by $\sim H \delta \phi / \dot{\phi}$. The solution to this problem has been discussed in [18]. Considering evolution during inflation[18] based on the δN calculation given in Ref.[22, 23], the correction from the evolution is found to be given by

$$\mathcal{R} \sim H \frac{\delta \phi}{\dot{\phi}}, \quad (2.25)$$

which is consistent with the expectation. Creation (evolution) of the curvature perturbations during inflation is considered by many authors[22, 24, 25, 26, 27].⁶ Besides the standard scenario in which the curvature perturbations are created at the horizon exit or during inflation, there is an alternative in which the curvature perturbations are created at the end of inflation[28, 29, 30, 31]. There are many other possibilities for the creation of the curvature perturbations. For example, creation of the curvature perturbations is possible at preheating[7, 32, 33, 34, 35], reheating[36, 37], or even at much later than the reheating period[17, 38, 39, 40]. They may be important for remote inflation, however for simplicity most of these alternatives are not discussed in this paper.

Below we first consider the creation of the curvature perturbations at the beginning of inflation, then we examine the inhomogeneous phase transition. Inhomogeneities of the diffusion rate $\delta \Gamma$ may lead to a significant creation of the curvature perturbations, however for the simplicity of the argument, they are not considered in this paper. The baryon number of the Universe caused by the decay of the sneutrino inflaton[5, 44] (leptogenesis) is discussed in section 2.3.

⁶ Here the “evolution” should be distinguished from the trivial k -dependence of $\mathcal{R}^{(ini)}$.

Due to the thermal effect, the amplitude of the inflaton perturbation is enhanced during strongly dissipating warm inflation. From the Langevin equation, the root-mean square fluctuation amplitude of the inflaton field $\delta\phi$ after the freeze out is obtained to be[1]

$$\delta\phi_{\Gamma>H} \sim (\Gamma H)^{1/4} T^{1/2} \sim r_\Gamma^{1/4} r_T^{1/2} H \quad (2.26)$$

where r_T denotes the ratio between T and H , defined by $r_T \equiv T/H$, which is assumed to be $r_T > 1$ due to the definition of warm inflation. We find for the sneutrino-remote scenario characterized by $\Gamma = A_N \phi_N$ and $V(\phi_N) \sim V_0 + M_N^2 \phi_N^2/2$;

$$r_T \simeq \frac{M_N \phi_N^{1/4}}{A_N^{1/4} H^{5/4}} \sim \eta^{1/2} \left(\frac{\phi_N}{A_N H} \right)^{1/4} \quad (2.27)$$

$$r_\Gamma \simeq \frac{A_N \phi_N}{3H}. \quad (2.28)$$

Therefore, for SDWI, the curvature perturbation created during warm inflation is given by

$$\mathcal{R}_{warm} \simeq H \frac{\delta\phi}{\dot{\phi}} \simeq \frac{r_\Gamma^{3/4}}{r_T^{3/2}} \sim \eta^{-3/4} A_N^{9/8} \left(\frac{\phi_N}{H} \right)^{3/8}. \quad (2.29)$$

From the equation for the number of e-foldings, it is very natural to expect $\phi_N^{(i)}/\phi_N^{(c)} \sim 1$ for $N_e \sim 60$.⁷ We thus find a more concrete estimation of the curvature perturbations given by

$$\mathcal{R}_{warm} \simeq A_N^{3/2} \left(\frac{r_{T_c}}{\eta} \right)^{3/2}. \quad (2.30)$$

Here T_c is the critical temperature defined for the remote sector, which determines the end of inflation. Imposing the normalization given by $\mathcal{R} \sim 10^{-5}$, we find an estimation of the coefficient given by

$$A_N \sim 10^{-3.3} \left(\frac{\eta}{r_{T_c}} \right). \quad (2.31)$$

Considering the $O(H)$ correction from the Kähler potential in supergravity, natural value of the η -parameter is $r_\Gamma > \eta \sim 1$. Low-scale inflation may lead to $\eta \gg 1$, which does not ruin the warm-inflationary scenario as far as the slow-roll conditions for warm inflation are satisfied.

⁷Note that for the present model the end of remote inflation is defined by $\phi_N^{(c)}$.

Contrary to the usual scenario of generating curvature perturbations at the end of inflation[28, 29, 31], here a model is considered in which the inhomogeneities may be created by the spatial inhomogeneities of the critical temperature (δT_c).⁸ The sources of the inhomogeneities are light fields (moduli), which determines the value of g and m_σ . From the definition of the critical temperature, we find

$$\frac{\delta T_c}{T_c} = \frac{\delta m_\sigma}{m_\sigma} - \frac{\delta g}{g}. \quad (2.32)$$

To find the curvature perturbations we must solve

$$\delta N_{end} \sim H \frac{\delta T_c}{\dot{T}} \quad (2.33)$$

at the end of inflation. In the present model, \dot{T} is determined by the inflaton motion $\dot{\phi}_N$ as

$$\dot{T} \sim \frac{M_N^4}{A_N H T^3} \dot{\phi} \sim \frac{\eta^2}{A_N r_{T_c}^3} \dot{\phi}. \quad (2.34)$$

Therefore, the curvature perturbation created at the inhomogeneous phase transition is estimated as

$$\begin{aligned} \mathcal{R}_{IP} \sim H \frac{\delta T_c}{\dot{T}} &\sim A_N r_{T_c}^4 \eta^{-2} \left(\frac{\delta m_\sigma}{m_\sigma} - \frac{\delta g}{g} \right) \frac{H^2}{\dot{\phi}} \\ &\sim A_N^2 \left(\frac{r_{T_c}^4}{\eta^3} \right) \left(\frac{\delta m_\sigma}{m_\sigma} - \frac{\delta g}{g} \right). \end{aligned} \quad (2.35)$$

The ratio between \mathcal{R}_{warm} and \mathcal{R}_{IP} is given by

$$r_{I/W} \equiv \frac{\mathcal{R}_{IP}}{\mathcal{R}_{warm}} \sim \eta^{-3/2} A_N^{1/2} r_{T_c}^{5/2} \left(\frac{\delta T_c}{T_c} \right). \quad (2.36)$$

The creation of the curvature perturbations by inhomogeneous phase transition requires $\mathcal{R}_{warm} < 10^{-5}$, which leads to a more concrete estimation of the ratio;

$$r_{I/W} \sim 10^{-5} \times A_N^{-1/2} \left(\frac{\delta T_c}{H} \right). \quad (2.37)$$

⁸Note that the creation of the curvature perturbations at the end of inflation is not general for usual sneutrino inflation, in which inflation is the chaotic-type. Instead, there would be an inhomogeneous preheating[32] at the end of chaotic-type inflation, which may cause significant generation of the curvature perturbations with non-Gaussian character.

2.2 Spectral index

Spectral index of the curvature perturbations in warm inflation has been calculated in Ref.[16] for general form of the dissipative coefficient. The calculation presented in the reference is straight and useful. According to the calculation in Ref.[16], the spectral index for hybrid-like inflation with $\Gamma \propto T^0$ and $V_{\phi_N} \propto \phi_N^{1+\epsilon_N}$ is given by

$$n_s - 1 \simeq -\frac{3\eta}{4r_\Gamma} \left(\frac{1 - 2\epsilon_N}{1 + \epsilon_N} \right), \quad (2.38)$$

which determines the ratio between η and r_Γ from the cosmological observations.

To summarize the results, we find the relation between A_N , r_{T_c} and η from \mathcal{R}_{warm} , and the equation for η , r_Γ from the spectral index. There are two equations for the four parameters A_N , T_c , V_0 and M_N , in which T_c and V_0 are the parameters of the remote sector. Non-Gaussianity is not significant when the curvature perturbations are generated by the standard mechanism. Therefore, observation of the non-linear parameter $f_{NL} \gg 1$ indicates that the curvature perturbations are not created at (and just after[18]) the horizon exit.

2.3 Leptogenesis

Leptogenesis can proceed through the out-of equilibrium decay of the sneutrino inflaton ϕ_N . Contrary to the usual sneutrino inflation model, reheating in remote inflation is caused by the remote-sector phase transition. Therefore, there is a dilution that leads to a suppression factor $\epsilon_D \sim M_N^2(\phi_N^{(e)})^2/V_0$ for the lepton number density. The lepton asymmetry of the Universe is measured by the lepton to entropy ratio that is given by

$$\frac{n_L}{s} \simeq |\epsilon_{CP}| \epsilon_D \frac{T_{RH}}{M_N}, \quad (2.39)$$

where ϵ_{CP} is the CP asymmetry generated at the decay of the sneutrino inflaton and T_{RH} is the reheating temperature. A typical bound for the CP-asymmetry parameter is given by[41]

$$|\epsilon_{CP}| \leq 2 \times 10^{-8} \frac{M_N}{10^8 GeV}, \quad (2.40)$$

which leads to the condition

$$\frac{n_L}{s} \leq 2 \times 10^{-8} \frac{T_{RH}}{10^8 GeV} \epsilon_D. \quad (2.41)$$

As one would easily understand from the typical potential of the thermal inflation model, the reheating temperature in sneutrino-remote inflation can be lowered to a very small scale, even smaller than the electroweak scale. Such a low-reheating temperature is very interesting, but it obviously ruins the successful scenario of leptogenesis caused by the sneutrino inflaton. Inflation and the creation of the curvature perturbations are natural in this sneutrino-remote model, which may work even if the reheating temperature is very low. This is an interesting result. However, leptogenesis may not be sufficient if the reheating temperature is very low. The simplest way out of this dilemma would be to find a more general inflaton candidate in MSSM fields.⁹ We consider this possibility in section 2.4. However, contrary to the sneutrino-remote inflation model that gives very simple results, arguments for MSSM-remote inflation with general inflaton potential are rather diverse. Moreover, the form of Γ may be different depending on the model parameters and the choice of the MSSM fields. Therefore, the purpose of the next paragraph is to show essential equations for MSSM-remote inflation when the diffusion parameter is given by the simplest form $\Gamma \sim A_M \phi_M$.

2.4 MSSM field as the inflaton

Applying remote inflation scenario to MSSM inflaton field, one finds immediately that significant difference may appear in the inflaton potential and ϵ_{CP} . The difference in the inflaton potential may lead to the difference in the slow-roll conditions, the curvature perturbations and the spectral index, while the difference in ϵ_{CP} may lead to a significant enhancement of the baryon number asymmetry for low-reheating temperature.

When the MSSM inflaton (ϕ_M) has the (quasi-)flat potential, the typical form of the effective inflaton potential for $T > T_c$ is given by

$$V(\phi_M) = V_0 \left\{ 1 + c'_0 \ln \frac{\phi_M}{M_*} \right\} \equiv V_0 + m^4 \lambda_0 \ln \frac{\phi_M}{M_*} \quad (n \equiv 0), \quad (2.42)$$

or considering non-renormalizable terms the effective potential may be characterized by

$$V(\phi_M) = V_0 \left\{ 1 + c_n \left(\frac{\phi_M}{M_*} \right)^n \right\} \equiv V_0 + \lambda_n \frac{\phi_M^n}{M_*^{n-4}} \quad (n \neq 0), \quad (2.43)$$

⁹There are many models of baryogenesis that may work with low-scale inflation. See for example Ref.[42] and the references therein.

where c'_0 , c_n and M_* are determined by the choice of the MSSM field, while V_0 is determined by the remote-sector potential. Although the potential with $n < 0$ is unlikely for the MSSM fields, it may appear in a hidden sector of dynamical supersymmetry breaking. In this paper we consider implicitly the potential with $n = 0$ and $n \geq 4$. The potential characterized by $n = 2$ is the same as the sneutrino-remote model, although the typical mass-scale is utterly different.

For simplicity, we consider only the typical form of the dissipative coefficient that is given by

$$\Gamma \equiv A_M \phi_M, \quad (2.44)$$

where the coefficient A_M is determined by the Yukawa couplings. The slow-roll velocity of the MSSM inflaton for the strongly dissipating regime ($r_\Gamma > 1$) is given by

$$\dot{\phi}_M \simeq -\frac{c'_0 V_0}{\Gamma \phi_M} \simeq -\frac{c'_0}{A_M} \frac{V_0}{\phi_M^2} \simeq -\frac{m^4 \lambda_0}{A_M \phi_M^2} \quad (n = 0) \quad (2.45)$$

$$\begin{aligned} \dot{\phi}_M &\simeq -\frac{nc_n V_0}{\Gamma \phi_M} \left(\frac{\phi_M}{M_*}\right)^n \simeq -\frac{nc_n V_0}{A_M \phi_M^2} \left(\frac{\phi_M}{M_*}\right)^n \\ &\simeq -\frac{n \lambda_n M_*^2}{A_M} \left(\frac{\phi_M}{M_*}\right)^{n-2} \quad (n \neq 0). \end{aligned} \quad (2.46)$$

Again, T -dependence of Γ leads to a condition for the radiation[1]

$$c \equiv \left| \frac{d \ln \Gamma}{d \ln T} \right| < 4, \quad (2.47)$$

which is satisfied **except** for $1 \leq n \leq 2$ [16].

In addition to the standard slow-roll conditions for the inflaton motion, slow-roll approximation for warm inflation requires slow-variation of Γ , which leads to the additional slow-roll parameter given by[1]

$$\beta \equiv M_p^2 \left(\frac{\Gamma_{\phi_N} V_{\phi_N}}{\Gamma V} \right) \simeq \frac{c'_0 M_p^2}{\phi_M^2} \sim \frac{m^4 \lambda_0}{H^2 \phi_M^2} < 1 + r_\Gamma \quad (n = 0), \quad (2.48)$$

$$\beta \equiv M_p^2 \left(\frac{\Gamma_{\phi_N} V_{\phi_N}}{\Gamma V} \right) \simeq \frac{nc_n M_p^2 \phi_M^{n-2}}{M_*^n} \sim \frac{n \lambda_n \phi_M^{n-2}}{H^2 M_*^{n-4}} < 1 + r_\Gamma \quad (n > 2), \quad (2.49)$$

which lead to non-trivial conditions for the inflaton field;

$$\phi_M > M_p \sqrt{\frac{c'_0}{(1+r_\Gamma)}} \sim \frac{m^2 \sqrt{\lambda_0}}{H \sqrt{1+r_\Gamma}} \equiv \phi_M^\beta \quad (n = 0) \quad (2.50)$$

$$\phi_M < \left(\frac{(1+r_\Gamma) M_*^n}{nc_n M_p^2} \right)^{\frac{1}{n-2}} \sim \left(\frac{(1+r_\Gamma) H^2 M_*^{n-4}}{n \lambda_n} \right)^{\frac{1}{n-2}} \equiv \phi_M^\beta \quad (n > 2) \quad (2.51)$$

The temperature during inflation is given by

$$T^4 \sim \frac{1}{H\pi^2 g_*} \frac{(c'_0)^2 V_0^2}{A_M \phi_M^3} \sim \frac{m^8 \lambda_0^2}{H\pi^2 g_* A_M \phi_M^3} \quad (n=0), \quad (2.52)$$

$$T^4 \sim \frac{1}{H\pi^2 g_*} \frac{n^2 c_n^2 V_0^2}{A_M \phi_M^3} \left(\frac{\phi_M}{M_*} \right)^{2n} \sim M_*^4 \frac{n^2 \lambda_n^2}{\pi^2 g_* A_M} \left(\frac{M_*}{H} \right) \left(\frac{\phi_M}{M_*} \right)^{2n-3} \quad (n>2). \quad (2.53)$$

Remote inflation is possible during $T > T_c$, which defines the critical point given by

$$\begin{aligned} \phi_M < \phi_M^{(c)} &\equiv \left(\frac{c'_0 V_0^2}{H\pi^2 g_* A_M T_c^4} \right)^{1/3} \\ &\sim \left(\frac{m^8 \lambda_0^2}{H\pi^2 g_* A_M T_c^4} \right)^{1/3} \quad (n=0), \end{aligned} \quad (2.54)$$

$$\begin{aligned} \phi_M > \phi_M^{(c)} &\equiv \left(\frac{H\pi^2 g_* A_M T_c^4 M_*^{2n}}{n^2 c_n^2 V_0^2} \right)^{\frac{1}{2n-3}} \\ &\sim \left(\frac{H\pi^2 g_* A_M T_c^4 M_*^{2n-8}}{n^2 \lambda_n^2} \right)^{\frac{1}{2n-3}} \quad (n>2). \end{aligned} \quad (2.55)$$

Here T_c is determined by the remote-sector potential.

In addition to these conditions, the conventional slow-roll conditions are the important conditions that determine the cosmological quantities created during inflation. The condition related to the ϵ -parameter is broken when $\epsilon > 1 + r_\Gamma$, which leads to the conditions for the slow-roll inflation given by

$$\phi_M > \phi_M^\epsilon \equiv M_p \frac{c'_0}{\sqrt{1+r_\Gamma}} \sim M_p \frac{\lambda_0}{\sqrt{1+r_\Gamma}} \left(\frac{m^4}{V_0} \right) \quad (n=0) \quad (2.56)$$

$$\phi_M < \phi_M^\epsilon \equiv M_* \left(\frac{\sqrt{1+r_\Gamma} M_*}{n c_n} \frac{M_*}{M_p} \right)^{\frac{1}{n-1}} \sim \left(\frac{\sqrt{1+r_\Gamma} V_0 M_*^{n-4}}{n \lambda_n M_p} \right)^{\frac{1}{n-1}} \quad (n>2) \quad (2.57)$$

For $n > 2$ and $\lambda_n > 0$, slow-roll inflation is possible during $\phi_M^{(c)} < \phi_M < \text{Min}\{\phi_M^\epsilon, \phi_M^\beta\}$, while for $n = 0$ and $\lambda_0 > 0$, we find that the condition is given by $\text{Max}\{\phi_M^\epsilon, \phi_M^\beta\} < \phi_M < \phi_M^{(c)}$. Remote inflation is not successful if these conditions are not satisfied. To understand the end-boundary of remote inflation, it would be useful to consider the difference between $n = 0$ and $n > 2$. In contrast with $n > 2$, inflation with the potential characterized by $n = 0$ may terminate with the violation of the slow-roll condition. Namely, for the positive coefficient $\lambda_0 > 0$, inflaton motion of the $n = 0$ potential is roll-in, and the slow-roll condition is violated at $\phi_M = \text{Max}\{\phi_M^\epsilon, \phi_M^\beta\}$. The temperature is increasing during inflation, which means that there is no phase transition during inflation. However, for

the negative coefficient $\lambda_0 < 0$, inflaton motion is roll-out, and the inflation ends with the phase transition at $\phi_M = \phi_M^{(c)}$. The temperature is decreasing during inflation.

To understand clearly these arguments, it would be useful to consider a simple and concrete example in which some parameters are fixed at typical values. To show an explicit example, we find for the simplest case with $n = 6$;

$$T^4 \simeq \frac{\lambda_6^2 \phi_M^9}{H A_M M_*^4}, \quad (2.58)$$

which leads to the critical point defined by the phase transition in the remote sector;

$$\phi_M^{(c)} \simeq (A_M H T_c^4 M_*^4 \lambda_6^{-2})^{1/9}, \quad (2.59)$$

and the critical points from the slow-roll conditions;

$$\phi_M^\epsilon \simeq \left(\frac{\sqrt{1+r_\Gamma} V_0 M_*^2}{\lambda_6 M_p} \right)^{1/5} \quad (2.60)$$

$$\phi_M^\beta \simeq \left(\frac{(1+r_\Gamma) H^2 M_*^2}{\lambda_6} \right)^{1/4}. \quad (2.61)$$

Obviously, remote inflation is possible for MSSM inflaton fields, during $\phi_M^{(c)} < \phi_M < \text{Min}\{\phi_M^\epsilon, \phi_M^\beta\}$. Using $\phi_M^{(c)}$ for the end of remote inflation, and assuming that $\phi_M^{(i)} \sim \phi_M^{(c)}$ for $N \sim 60$, we find the curvature perturbations given by

$$\mathcal{R}_{MSSM} \simeq \frac{r_\Gamma^{3/4}}{r_T^{3/2}} \sim r_{T_c}^{-7/6} A_M^{5/6} \lambda_6^{-1/6} \left(\frac{M_*}{H} \right)^{1/3}. \quad (2.62)$$

An intuitive argument would be useful for understanding the discrepancy between the usual MSSM inflation and MSSM-remote inflation. Standard MSSM inflation[6] realized with tiny Hubble parameter requires tiny slow-roll parameter for the creation of the curvature perturbations. Namely, from the normalization of the curvature perturbations and the standard formula

$$\mathcal{R} \sim H \frac{\delta\phi}{\dot{\phi}}, \quad (2.63)$$

it is found that the slow-roll parameter $\epsilon \sim 10^{-22} (H/10^2 \text{GeV})^2$ is required to create the cosmological perturbations.¹⁰ Since the condition required for the creation of the curvature perturbations ($\epsilon < 10^{-22}$) is much more severe than the slow-roll condition

¹⁰ Here the creation of the curvature perturbations at the horizon exit is discussed. Curvatons or inhomogeneous preheating may relax the fine-tuning problem in low-scale inflation scenario [33, 40].

$\epsilon < 1$, there is a dilemma that ϵ at the creation of the curvature perturbations must be far different from the value at the end of inflation. Namely, rapid variation of ϵ is required for the last e-foldings $N_e \sim 60$, while the potential must be very flat at the horizon exit. In contrast with supercooled inflation, strong dissipation in warm inflation may cause strong damping of the inflaton motion, even if the Hubble-induced mass leads to $\eta \sim 1$. Moreover, since the end of remote inflation is not defined by $\epsilon \sim 1$, there is no dilemma related to the variation of ϵ . If the dissipation rate is given by $\Gamma \sim A\phi_M$, it enhances the curvature perturbation by the huge factor $\Gamma/H \sim r_\Gamma \sim A\phi_M/H \gg 1$. This factor may cure the fine-tuning problem of the supercooled MSSM inflation model.

2.5 Baryogenesis for MSSM inflaton

Compared with sneutrino-remote inflation, a significant discrepancy may appear in the estimation of the baryon number asymmetry, which is caused by the magnitude of ϵ_{CP} and the mass of the inflaton field. A rough estimation of the baryon asymmetry of the Universe measured by the baryon to entropy ratio is given by

$$\frac{n_B}{s} \sim |\epsilon_{CP}| \epsilon_D \frac{T_{RH}}{m_{soft}}, \quad (2.64)$$

where T_{RH} can be as small as the scale of the nucleosynthesis[3]. The MSSM direction may have either baryon or lepton number. For high T_R , leptogenesis may work with active sphalerons. The baryon number is conserved if T_R is very low, where there is no active sphaleron after reheating.

3 Conclusions

In this paper we applied a new scenario of hybrid-like inflation to sneutrino and MSSM inflaton fields. In remote inflation, radiation raised continuously by a dissipating inflaton field keeps symmetry restoration in a remote sector, and the false-vacuum energy of the remote sector dominates the energy density during inflation. The situation is similar to hybrid inflation, but the hybrid-type potential is not required in this scenario. Remote inflation ends when the temperature reaches the critical temperature, or when the slow-roll condition is violated. Using the scenario of remote inflation, we considered sneutrino

and MSSM fields as the inflaton candidate, and found essential conditions for the inflaton potential and the remote sector. The results show that remote inflation may naturally work with sneutrino or MSSM fields. More details of the scenario (i.e., more details of the required parameter space for the curvature perturbations and baryogenesis) must be studied using numerical calculations.

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A Thermal Inflation and reheating in the remote sector

A generic supersymmetric gauge theory will have a large number of directions which is flat before supersymmetry breaking. The minimal supersymmetric standard model (MSSM) will have the flat directions, but it does not lead to $V_0 \neq 0$ during remote inflation. We thus considered for the remote sector a flat direction (flaton) which is not the MSSM flat direction but can be thermalized during remote inflation.¹¹ The most obvious example of this kind has already been discussed for thermal inflation[3]. A field with a nonzero VEV is defined either by a Higgs field (which is not necessarily the standard-model Higgs or the conventional-GUT Higgs field) or a gauge singlet. The discussions for the original work of thermal inflation were focused on the latter case, but it was also commented that in some GUT models a Higgs field may play the role. In particular, gauge-mediated models of supersymmetry breaking has a hidden sector which interacts with the standard-model fields with renormalizable interactions. The crucial difference between remote-sector inflation and thermal inflation appears in the source of the radiation. In thermal inflation, there is no source of the radiation during inflation, while in remote inflation dissipation in the inflaton sector sources radiation.

¹¹Note that the remote field must be distinguished from the inflaton field. The remote field plays the role of the trigger field that appears in the conventional hybrid inflation.

With regard to reheating in the remote sector, our scenario is very similar to thermal inflation, which is **not similar to the reheating mechanism in MSSM inflation**.¹² Note that in MSSM-remote inflation, MSSM field appears as the inflaton in the inflaton sector, however reheating is caused by the remote-trigger field in the remote sector, which can be identified with the flaton in the conventional thermal inflation scenario. In the remote inflation scenario, reheating is not due to the decay of the MSSM inflaton, but induced by the remote-trigger field in the remote sector. We think there is no confusion on this point, since the scenario is comparable to the well-known scenario of hybrid inflation.

A flaton particle corresponding to oscillations around a VEV of M_0 will couple only weakly to particles with mass less than M_0 . The simplest choice of the decay rate is $\Gamma \sim m^3/M_0^2$, where m is the flaton mass near the VEV. For example, an effective interaction $\lambda|\sigma|^2|\chi|^2$ between the remote field σ and a spin-zero particle χ with the mass m_χ will lead to the decay rate[3]

$$\Gamma \simeq \frac{\lambda^2}{8\pi} \left(\frac{M}{m}\right)^2 m \sqrt{1 - \frac{4m_\chi^2}{m^2}}. \quad (\text{A.1})$$

Maximizing the decay rate with the obvious condition for the decay channel¹³

$$\lambda \leq \frac{1}{2} \frac{m_\chi^2}{M^2}, \quad (\text{A.2})$$

it is found that $\Gamma \leq 10^{-4}m^3/M^2$ [3]. Note that there are many choices for the remote sector, which are highly model-dependent.

The discussions for the flaton reheat temperature in the remote sector is precisely the same as what has been discussed for thermal inflation. The typical example that has been discussed in the original work of thermal inflation leads to the reheat temperature

$$T_D \sim \left(\frac{10^{11}\text{GeV}}{M_0}\right) \left(\frac{m}{300\text{GeV}}\right)^{3/2} \text{GeV}. \quad (\text{A.3})$$

Of course, there are many choices for the flaton field that lead to many different values of the reheat temperature. It is not important to show the catalog of the flaton candidates in this appendix, since the purpose of this paper is to show a significant possibility of MSSM-remote inflation.

¹²Note that the oscillation after thermal inflation does not lead to preheating. Some details are discussed for MSSM fields in Ref.[43]. See also Ref.[44].

¹³This condition must not be applied to the interaction with thermal background during remote inflation.

B On thermal corrections to the MSSM inflaton

If part of the standard-model gauge group (or the extended gauge group in some extended models of the standard model) remains unbroken by the VEV of flat fields, the effective mass of the associated gauge fields (and gauginos) will remain massless. The background thermal bath with a temperature T will affect the flat direction dynamics. A detailed study of such thermal corrections has been presented in Ref.[44] with regard to the scenario in which MSSM directions are identified with curvatons. For the MSSM-remote inflation scenario, the situation is precisely the same as what happens in the conventional warm-inflation scenario. Namely, for the MSSM-inflaton VEV ϕ that induces a mass $y\phi$ to the field that couples to the MSSM-inflaton, the effective potential that arises from the thermal background are

$$\begin{aligned}\Delta V_{th} &\sim y^2 T^2 |\phi|^2, \quad (y\phi \leq T) \\ \Delta V_{th} &\sim \pm \alpha T^4 \ln(|\phi|^2), \quad (y\phi > T)\end{aligned}\tag{B.1}$$

where α is a gauge fine structure constant. The obvious condition, which arises for any kind of warm inflation scenario, is given by $y\phi \gg T$. Of course, this condition must be satisfied during MSSM-remote inflation, although it frequently leads to a trivial condition that can be disregarded compared with other conditions[1].

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